Population III Gamma Ray Bursts

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ABSTRACT

We discuss a model of Poynting-dominated gamma-ray bursts from the collapse of very massive first generation (pop. III) stars. From redshifts of order 20, the resulting relativistic jets would radiate in the hard X-ray range around 50 keV and above, followed after roughly a day by an external shock component peaking around a few keV. On the same timescales an inverse Compton component around 75 GeV may be expected, as well as a possible infra-red flash. The fluences of these components would be above the threshold for detectors such as Swift and Fermi, providing potentially valuable information on the formation and properties of what may be the first luminous objects and their black holes in the high redshift Universe.

Subject headings: gamma-rays: bursts — cosmology — stars: population III — jets: magnetized – radiation mechanisms: non-thermal

1. Introduction

Population III stars are widely considered to consist mainly of 'very massive stars' (VMS) in the the range of hundreds of solar masses (Ohkubo et al.; Yoshida et al. 2006). The VMS are expected to be very fast rotating, close to the break-up speed, and accretion leads to a mass upper limit which may be around $10^3 M_\odot$. Those in the $140 M_\odot \lesssim M_* \lesssim 260 M_\odot$ range are expected to be subject to pair instability and explode as supernovae without leaving any compact remnant behind, while those above $\sim 260 M_\odot$ are expected to undergo a core collapse leading directly to a central black

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hole (Heger and Woosley 2002), whose mass would itself be hundreds of stellar masses. Accretion onto such massive black holes could lead to a scaled-up collapsar gamma-ray burst (Heger et al. 2003; Komissarov and Barkov 2009). In this paper we discuss a specific scenario for pop. III VMS collapsars at redshifts of order $z \sim 20$, resulting in Poynting dominated relativistic jets which produce GRBs with characteristic radiation properties extending from soft X-rays to multi-GeV energies.

2. A population III collapsar model

We consider a pop. III star undergoing core collapse at a redshift $z \sim 20$, which leaves behind a black hole of mass M_h surrounded by an accretion disk or torus of mass M_d. Knowledge about the progenitor structure and the collapse history is rather limited, an approximate but representative scenario having been outlined in Komissarov and Barkov (2009). This assumes a nominal VMS of mass $M_* = 10^3 M_3 M_\odot$ and radius $R_* = 10^{12} R_{*,12}$ cm rotating at half the break-up speed, which results in a disk of outer radius of R_d, disk mass M_d and central black hole of mass M_h. Given the uncertainties, for the purposes of estimates we can assume that the typical disk mass and the black hole mass are of the same order as the progenitor mass, $M_d \simeq M_h \simeq M_* \simeq 10^3 M_3 M_\odot$.

For such large BH masses the accretion torus density and temperature are too low for neutrino cooling to be important, and the accretion regime can be described through an advection dominate (ADAF) model, e.g. Narayan and Yi (1994), in which radiation pressure is dominant. The properties of the precursor VMS, and the flow dynamics after the collapse, are plainly uncertain; it is nonetheless helpful to parametrize the key numbers and scaling relations in terms of an undoubtedly oversimplified but specific model. For a VMS rotating at, say, half the break-up speed the disk outer radius will be $R_d = R_*/4$, and for a disk viscosity parameter $\alpha = 10^{-1}\alpha_{-1}$ the accretion time $t_d \simeq (14/9\alpha)(R^3/GM)^{1/2}$ is approximately

$$t_{ac} \simeq 5 \times 10^3 \alpha_{-1}^{-1} R_*^{3/2} M_3^{-1/2} \text{ s,} \tag{1}$$

and the mean accretion rate is $\dot{M} \simeq \frac{M_d}{t_{ac}} \simeq 0.2\alpha_{-1} R_{*,12}^{-3/2} M_3^{3/2} \ M_\odot \ \text{s}^{-1}$. The disk inner radius at the marginally bound orbit of a rotating Kerr black hole with rotation parameter a is

$$r_\ell = (R_s/2) f_1(a) \simeq 3 \times 10^8 M_3 \text{ cm,} \tag{2}$$

where $R_s = (2GM_h/c^2)$ and $f_1(a) = [2 - a + 2(1-a)^{1/2}] \simeq 2.1$ is estimated for $a \simeq 0.8$. The inefficient neutrino cooling is insufficient to power a strong jet, but strong magnetic field build-up in the torus could lead to much stronger MHD jets. The disk mass density ρ

and gas pressure P in the ADAF regime provide an estimate of the disk poloidal magnetic field as $B^2 = (8\pi P/\beta)$, where $\beta = 10\beta_1$ is the magnetization parameter, leading to a Blandford-Znajek type (Blandford and Znajek 1977) Poynting jet luminosity

$$L = \frac{\pi c^2}{48\beta} f_1^{3/2} f_2^2 \frac{G^{1/2} M_h^{3/2}}{R_*^{3/2}} \simeq 5 \times 10^{52} \beta_1^{-1} R_{*,12}^{-3/2} M_3^{3/2} \text{ erg s}^{-1}, \tag{3}$$

where $f_2(a) = (a/2)(1 + \sqrt{1 - a^2})$ and for $0.5 \lesssim a \lesssim 1$ the product $f_1^{3/2} f_2^2 \simeq 1/4$ (Komissarov and Barkov 2009). For a jet solid angle $\Omega = 10^{-2}\Omega_{-2}$, as may be expected in a large star, and assuming that a fraction $\eta = 10^{-1}\eta_{-1}$ of the luminosity falls in the X-ray band (see §5), from a redshift $z \sim 20$ and using current cosmological parameters, using $L_{iso} = L(4\pi/\Omega)$ a flux $F = \eta L_{iso}/4\pi r_L^2 \sim 10^{-6}\eta_{-1}\beta_1^{-1}\Omega_{-2}^{-1}R_{*,12}^{-3/2}M_3^{3/2}$ erg cm^{-2} s^{-1} is expected, more than an order of magnitude above the Swift BAT sensitivity. In the next section we show that such objects could indeed have such X-ray luminosities, as well as other components at higher energies.

3.   Population III Poynting dominated GRBs

Taking the magnetic luminosity of eq. (3) as representative for a pop. III collapsar at a redshift $z \simeq 20$, we use this as the central engine underlying a scaled-up version of a Poynting-dominated GRB model discussed in Mészáros and Rees, 1997; henceforth MR97). That is, we assume that a purely MHD (Poynting dominated) jet emanates from the central black hole plus accretion torus system, which is initially devoid of baryons. At the base of the jet r_ℓ given by eq. (2) the initial jet bulk Lorentz factor is parametrized as $\Gamma_\ell \simeq L/L_w$, where L is the magnetic luminosity (3) and $L_w \equiv L_{e^\pm,\gamma}$ is the associated pair and photon luminosity (allowing for the possibility that $\Gamma_\ell \geq 1$, as in pulsar wind models). Depending on uncertain details at the base of the jet, the initial Γ_ℓ could differ from unity, $\Gamma_\ell \gtrsim 1$. The jet magnetic luminosity L is related to the comoving transverse magnetic field B'_ℓ at the base of the jet through $L \simeq 4\pi r_\ell^2 c(B'^2/8\pi)\Gamma_\ell^2$, where $L = 5 \times 10^{52} L_{52.7}$ erg s$^{-1} \equiv 5 \times 10^{52} \beta_1^{-1} R_{*,12}^{-3/2} M_3^{3/2}$ erg s^{-1} can be taken as the isotropic equivalent luminosity for an observed within an angle $1/\Gamma$ of a jet axis whose final Lorentz factor $\Gamma \gg 1$.

We review here the dynamics of baryon-free Poynting jets for the VMS collapsar case, since the numerical values differ from those of the normal stellar collapsar case. Near the base r_ℓ of the outflow the jet will generally become loaded with pairs which are in near thermal equilibrium with photons. At the base of the jet the transverse comoving magnetic field strength, the comoving pair temperature and the comoving pair density are $B'_\ell \simeq 6 \times 10^{12} L_{52.7}^{1/2} r_{\ell,8.5}^{-1} \Gamma_\ell^{-1}$ G, $\quad T'_\ell \simeq 3.7 \times 10^9 L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{-3/4}$ K

and $n'_\ell \simeq 1.5 \times 10^{30} L_{52.7}^{3/4} r_{\ell,8.5}^{-3/2} \Gamma_\ell^{-9/4}$ cm^{-3}, where $T'_\ell \simeq (L_w/4\pi c r_\ell^2 \Gamma_\ell^2 a_B)^{1/4}$ and $n'_\ell \simeq (0.4 a_B k T'^4_\ell/kT'_\ell)$, with a_B is the Stefan-Boltzmann constant, $r_{\ell,8.5} = (r_\ell/3 \times 10^8 M_3$ cm$)$ is given by eq. (2) and $n' \sim n'_\pm \sim n'_\gamma$ with $T'_\pm \simeq T'_\gamma$. The optically thick pairs and photons will be frozen-in with the magnetic field, the whole behaving as a relativistic fluid. The magnetic geometry is that of an aligned rotator, the field spiraling out along the axis, and the pair plasma is frozen-in. Assuming for simplicity a constant opening angle jet with purely radial motion, the lab frame transverse magnetic field strength varies as $B \propto r^{-1}$, and the comoving transverse magnetic field strength varies as $B' \propto B/\Gamma$. Along the jet the transverse field lines do not change polarity, which is less likely to lead to reconnection, especially for an essentially baryon-free outflow such as we assume here. This differs from models, e.g. Drenkhahn and Spruit (2002), which assume a baryon load and where reconnection plays a role in the dynamics of the expansion. In the absence of reconnection, the magnetically dominated comoving energy density $\varepsilon' \propto r^{-2}\Gamma^{-2}$, while the comoving pair density $n' \propto T'^3 \propto r^{-3}$. The fluid has a bulk Lorentz factor Γ in the lab frame and an internal (comoving) Lorentz factor $\gamma' \sim \varepsilon'/n'$, and from energy and entropy conservation, Γ grows at the expense of γ', so $\Gamma \propto 1/\gamma' \propto n'/\varepsilon' \propto r$, for a jet of constant opening angle. Hence, initially $\Gamma \simeq \Gamma_\ell(r/r_\ell)$, and the comoving pair equilibrium temperature varies as $T' \propto r^{-1}$.

When the comoving temperature drops below $m_e c^2$ the pairs start to recombine, but the linear acceleration continues beyond this point, until the comoving pair Thomson optical depth τ'_T has become less than unity, which occurs when the comoving temperature $T' \propto r^{-1}$ has dropped to $kT'_a \simeq 0.04 m_e c^2 \simeq 17$ keV (e.g. Shemi and Piran, 1990). This occurs in our case at a radius r_a where $(r_a/r_\ell) = (T'_\ell/T'_a) = (\Gamma_a/\Gamma_\ell) \simeq 20 L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{-3/4}$, or

$$
\begin{aligned}
r_a &\simeq 6 \times 10^9 L_{52.7}^{1/4} r_{\ell,8.5}^{1/2} \Gamma_\ell^{-3/4}, \\
\Gamma_a &\simeq 2 \times 10^1 L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{1/4}.
\end{aligned}
\tag{4}
$$

At this radius most of the pairs have already recombined, and the gas density consist mainly of photons, $n'_a \simeq n'_\ell(r_a/r_\ell)^{-3} \simeq 2 \times 10^{26}$ cm^{-3}, while the remaining pair comoving density, from Saha's equation, is $n'_{\pm,a} \simeq 5 \times 10^{15} \Gamma_\ell r_{\ell,8.5}^{-1}$ cm^{-3}. The photon to pair ratio at this radius is $\sim 4 \times 10^{10}$, and pair annihilation practically ceases beyond this, so the maximum theoretical (inertial limit) Lorentz factor is $\Gamma_{in} \sim (n'_\gamma/n'_\pm)(kT'_a/m_e c^2) \simeq 1.6 \times 10^9$. Other effects, however, can set in before that, resulting in a lower terminal value.

Above the annihilation radius r_a, the comoving density of pairs $n' \equiv n'_\pm$ providing the inertia has been drastically reduced, but the magnetic pressure or the comoving energy density $\varepsilon' \propto B'^2$ continues acting continuously. In the lab frame the transverse

component of the field varies as $B \propto r^{-1}$. Since $\Gamma \propto n'/\varepsilon' \propto n'/B'^2 \propto B^2/n'$, the drop in n' implies that the gas must accelerate faster than the previous behavior of $\Gamma \propto r$. The pair density is above the minimum below wich the plasma can't carry the MHD currents even if the stream velocity of the positive relative to negative charges is of order c (Goldreich and Juian 1969), which is essentially the same for an e^+e^- or an e^-p^+ plasma (note that the Blandford-Znajek mechanism exploits the analogy with the aligned pulsar case); thus the MHD regime remains valid.

In the next three paragraphs we outline a possible scenario for the dynamics above the photospheric radius of a Poynting jet, which can lead to very high bulk Lorentz factors Γ. This high-Γ scenario is sensitive to model details, and must be considered speculative. For instance, such high Lorentz factors could lead to gradients which invalidate the transverse field assumption; also, a small amount of entrained baryons could end up dominating the inertia. For simplicity we neglect such potential complications, noting that our final observational predictions are essentially idependendent of what happens in this high-Γ regime. Thus, we consider a highly idealized picture of what happens above the photospheric radius r_a. While most of the photons escape freely, the pairs continue being scattered repeatedly by the much more numerous photons, and continue experiencing a drag for some distance $r > r_a$. It is useful to compute this in a frame $\Gamma_i \sim (r/r_\ell)\Gamma_\ell$ where the photons are isotropic. In this frame the electrons, which are essentially cold in the comoving frame, are boosted to a Lorentz factor $\gamma \sim \Gamma/\Gamma_i$, and the drag time is $t_{dr,i} \sim m_e c^2/(u_{ph,i}\sigma_T c\gamma) = (m_e c^2 4\pi r^2 \Gamma_i^3/L\sigma_T\Gamma)$, where u_{ph} is the radiation energy density. In the lab frame the drag time is Γ_i times longer, and the acceleration rate is obtained by setting the ratio of the lab frame Compton drag time $t_{dr} = (m_e c^2 4\pi r^2/L\sigma_T\Gamma)(r/r_\ell)^4$ and the lab frame expansion time r/c equal to the ratio of the kinetic and the Poynting flux, $n'm_e c^2\Gamma^2/[(B_\ell^2/4\pi)(r_\ell/r)^2]$. The drag time is shorter than the annihilation time at $r \geq r_a$ so the remaining pairs are frozen in, and for a dominantly transverse magnetic field, the ratio of the drag to expansion time is $\propto r^5/\Gamma$ while the ratio of the kinetic to Poynting flux is $\propto \Gamma$, so for $r \geq r_a$ the flow accelerates as $\Gamma \propto r^{5/2}$.

As the Lorentz factor continues increasing beyond r_a the annihilation photons, whose isotropic frame energy is $0.12 m_e c^2(r_a/r)$, eventually are blueshifted in the jet frame to $\gtrsim m_e c^2$, and their directions are randomized by scattering. Given that the compactness parameter is large, this results in copious pair production $\gamma\gamma \to e^+e^-$. This reduces the drag while increasing the inertia, leading to a mitigation of the acceleration rate beyond a radius r_p where the Lorentz factor is Γ_p,

$$\begin{aligned}
r_p &\simeq 4.6 \times 10^{11} L_{52.7}^{1/4} r_{\ell,8.5}^{1/2} \Gamma_\ell^{-3/4} \quad \text{cm}, \\
\Gamma_p &\simeq 1.5 \times 10^6 L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{1/4}.
\end{aligned} \tag{5}$$

Beyond r_p, pair formation will be self-limiting, resulting in a scattering optical depth $\tau_\pm \sim 1$, the threshold condition being $0.12 m_e c^2 (r_a/r)(\Gamma/\Gamma_i) \sim m_e c^2$, and since $\Gamma_i \propto r$ this implies at $r > r_p$ an acceleration rate $\Gamma \propto r^2$.

A basic uncertainty is that in addition to the jet's own annihilation photons there will be other photon sources which may also provide drag on the pairs. One such additional photon source is the accretion torus, but this radiation will be absorbed by the rest of the stellar envelope through which the jet is making its way, unless all of it is incorporated in the torus or is blown away. Also, the radius r_p is near the outer envelope of a VMS of radius $R_* \sim 10^{12}$ cm, which is a source of UV and soft X-ray photons. The drag on these external photons is likely to limit the Lorentz factor to values not much above Γ_p.

Irrespective of the final Lorentz reached in the previous high-Γ phase, outside the stellar radius R_* the jet is expected to propagate through a stellar wind, whose details are poorly known for pop. III objects, and further out the jet will encounter the interstellar medium of the minihalo or protogalaxy hosting the VMS. The jet will shock and sweep up the external medium, pushing it ahead of itself. While the jet continues to be fed by accretion, over a time of order t_{ac} given by eq. (1), after an initial brief transient the shocked jet head will continue to advance at a decelerating pace into the external medium. The Lorentz factor of the jet head is determined by momentum balance across the shock front. The shock has a Lorentz factor $\sim \Gamma$ in the lab, and in the shock frame the kinetic pressure (magnetic, or radiation) $p' \sim \varepsilon' \propto L/(r^2 \Gamma^2)$ on one side must balance the ram pressure from the external matter on the other side, which in the shock frame is $\rho_{ext} \Gamma^2$. Their equality defines the bulk Lorentz factor of the shock, $\Gamma \propto r^{-1/2}$, which is now decelerating. This continues as long as the jet continues being fed at constant luminosity, for $t < t_{ac}$. When $t = t_{ac}$ is reached, independently of the initial Lorentz factor value, the jet will have decelerated to a value Γ_d, reached at a deceleration radius r_d given by

$$
\begin{aligned}
r_d &\simeq 1.2 \times 10^{18} L_{52.7}^{1/4} t_{a,3.5}^{1/2} n_2^{-1/4} \quad \text{cm}, \\
\Gamma_d &\simeq 1.3 \times 10^2 L_{52.7}^{1/8} t_{a,3.5}^{-1/4} n_2^{-1/8} .
\end{aligned}
\tag{6}
$$

At this time $t_{ac} \sim r_d/c\Gamma_d^2$ in the source frame, after feeding of the jet ceases (at this radius, the frozen-in pair density in the jet still exceeds the nominal Goldreich-Julian critical density, so the MHD assumption would remain valid). We have assumed a typical $z \sim 20$ minihalo or protogalaxy gas density of $n_{ext} \sim 100$ cm^{-3}, e.g. Madau and Rees (2001). At the radius r_d this density is of order of or higher than that of a possible $\lesssim 10^{-3} M_\odot/\text{yr}$ stellar wind. Beyond this deceleration radius, for a constant external density the jet decelerates in the energy conserving regime (e.g. Blandford and McKee, 1976) at the self-similar rate $\Gamma \propto r^{-3/2}$.

4. Radiation properties

Annihilation photons .– The annihilation photons escaping from the pair photosphere at r_a given by eq. (4) appear, in the observer frame, with a peak energy of

$$E_{an}^{ob} \simeq \frac{\Gamma_a 3kT_a'}{(1+z)} \simeq 50 \text{ keV} L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{1/4} (20/[1+z]) \quad \text{(annihilation)}. \quad (7)$$

Inverse Compton scattering in such photospheres will generally lead also to a high energy power law extending as $N(E) \propto E^{-2}$ above the peak (Pe'er et al. 2006). In our case, however, there will also be up-scattering of annihilation photons in the drag region $r_a \lesssim r \lesssim r_p$. This will depend on the scattering optical depth, which is of order unity just below r_a and is very small above r_a, but increases to $\tau_\pm \sim 1$ at $r \gtrsim r_p$ where pair formation sets in. One can expect a significant component of upscattered photons from the two radii where $\tau_\pm \sim 1$, namely from r_a at energies $0.12\Gamma_a m_e c^2/(1+z)$, and from r_p at energies $\sim \Gamma_p m_e c^2/(1+z)$ in the observer frame,

$$\begin{aligned} E_{an,sc,a}^{ob} &\sim \quad 50 \text{ keV } L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{1/4} (20/[1+z]) \quad \text{(from } r_a) \\ E_{an,sc,p}^{ob} &\sim \quad 25 \text{ GeV } L_{52.7}^{1/4} r_{\ell,8.5}^{-1/2} \Gamma_\ell^{1/4} (20/[1+z]) \quad \text{(from) r}_p). \end{aligned} \quad (8)$$

These components would appear as two humps at these energies, and would have a comparable energy, which is a significant fraction of the jet energy. Above r_p, if the jet continues to accelerate for a while before slowing down, both the jet annihilation photons and external photons from the stellar envelope lead to pair formation, and some fraction of the jet energy could conceivably continue going into high energy photons, but the dynamics is dependent on the stellar model's dynamical behavior and photon input during the collapse and the jet propagation. A theoretical upper limit for the photon energy would be $\lesssim \Gamma_i m_e c^2$, where Γ_i is the inertial limit $\sim 10^9$ discussed below eq. (4), which however is unlikely to be reached, either because the field develops a longitudinal component, or because the flow may acquire some baryons (and associated electrons) whose inertia eventually becomes important.

Interaction with external photons.– Pairs and photons in the jet can also interact with external photons from the progenitor star or the accretion disk can exert an additional drag as well as resulting in additional spectral components. The latter could contribute significantly to the total spectrum if the jet Lorentz factor reaches values of at least $\sim \Gamma_p \sim 10^6$. This is because the maximum boost in photon energy from the interaction is of order $\lesssim \Gamma^2$, while the stellar or disk photon luminosity may be of order the photon Eddington value $L_{Ed,*} \sim 10^{41} M_3$ erg s^{-1}, giving a component of luminosity $L_{ext} \lesssim L_{Ed}\Gamma^2 \lesssim 10^{53}$ erg s$^{-1} \lesssim L$ which could reach a substantial fraction of the jet Poynting luminosity L.

The process is complicated by the fact that the external (stellar or disk) radiation field will be inhomogeneous across the jet cross section, as well as depending on height. With the nominal values used here, for a jet opening angle $\theta_j \sim 10^{-1}$ the lab-frame transverse Thomson optical depth of frozen-in jet pairs is $\tau_{T,\perp} \lesssim 1$ at $r \sim r_p$, so a drag (and boost) may apply over the entire jet cross section. However pair formation with external photons (and annihilation photons) above r_p could introduce significant angle dependent optical depth effects. Neglecting such inhomogeneities, a discussion of an upscattered component and a simple example of a one-zone pair cascade spectral component at r_p were discussed in MR97 for a Poynting jet from normal (pop. II) stars. Based on the expressions in that paper, the same processes would result, at r_p in the present physical system, in a component from upscattered stellar photons at $E_{sc}^{ob} \sim 250$ TeV $L_{52.7}^{1/2} r_{\ell,8.5}^{-1} \Gamma_\ell^{1/2}(20/[1 + z])$, which would be absorbed by $\gamma\gamma$ interactions against intergalactic IR photons; and in a pair cascade component emerging at $E_{casc}^{ob} \sim$ 2.5 keV $L_{52.7}^{5/4} r_{\ell,8.5}^{-5/2} \Gamma_\ell^{1/4}(20/[1 + z])$, whose luminosity would be given by a Γ_p^2 boost. However, transverse inhomogeneity effects as well uncertainties concerning the height dependence require detailed (and model dependent) calculations, making it difficult to say anything beyond the above semi-quantitative comments.

Internal dissipation radiation.– In this model we do not expect radiation from internal shocks because in a magnetically dominated outflow these do not arise (just as they do not play a significant role in the Crab wind); and if there aren't reversals in the field there shouldn't be internal dissipation (§3).

External blast wave radiation.– The contribution from this component is subject to fewer uncertainties than the previous ones. The forward shock from the eject plowing into the external medium produces a luminosity peaking at the deceleration radius r_d where the shock Lorentz factor is $\sim \Gamma_d \sim 130$ (eq.[6]). Following the usual treatment of external shocks, we estimate that for an external density in the VMS host environment of $n_{ext} \sim 10^2 n_2$ cm^{-3} the shocked external gas may build up turbulent magnetic fields to some fraction ϵ_B of the equipartition value with the post-shock thermal energy, resulting in a comoving field in the shocked gas of $B' \sim (\epsilon_B 8\pi n_{ext}\Gamma_d^2 m_p c^2)^{1/2} \sim 8\times10^1 \epsilon_{B,-1}^{1/2} n_2^{1/2}\Gamma_{d,2.1}$ G. The comoving random Lorentz factors of the electrons in the shocked gas peak have a minimum (peak) Lorentz factor $\gamma_m' \sim \epsilon_e\Gamma_d(m_p/m_e) \sim 2.4 \times 10^4 \epsilon_{e,-1}\Gamma_{d,2.1}$, with a power law $N(\gamma') \propto \gamma'^{-p}$ above that due to Fermi acceleration. The comoving synchrotron peak will be at an energy $x_s'y \sim (3/4)x_B'\gamma_m'^2 \sim 8 \times 10^{-4}\epsilon_{B,-1}^{1/2} n_2^{1/2}\Gamma_{d,2.1}^3\epsilon_{e,-1}^2$, which in the lab frame is $x_{sy} \sim 10^{-1}\epsilon_{B,-1}^{1/2} n_2^{1/2}\Gamma_{d,2.1}^4\epsilon_{e,-1}^2 \sim 50$ keV, where $\Gamma_{d,2.1} \equiv L_{52.7}^{1/8} t_{a,3.5}^{-1/2} n_2^{-1/8}$ (eq.[6]). Thus, in the observer frame this is

$$E_{sy}^{ob} \sim 2.5 \text{ keV } \epsilon_{B,-1}^{1/2}\epsilon_{e,-1}^2 L_{52.7}^{1/2} t_{a,3.5}^{-1}(20/[1 + z]) \quad \text{(synchrotron)}, \qquad (9)$$

independent of the assumed external (host) density. This synchrotron peak photon

energy is accidentally, similar to that for the external cascade at a nominal radius r_p. As mentioned there, however, the cascade photon energy will be smeared by integration over r, whereas the external shock synchrotron energy (9) is generally fairly well defined, as observations of normal GRB indicate. The luminosity of this blast wave synchrotron component would be a substantial fraction of the Poynting luminosity (eq.[3]), and it would peak at the deceleration time $t_d \sim t_a$ (eqs.[6], [1]) redshifted to the observer frame,

$$t_{sy,pk}^{ob} \sim 10^5 \alpha_{-1}^{-1} R_*^{3/2} M_3^{-1/2} ([1+z]/20) \ \text{s.} \tag{10}$$

This would also be the order of the time delay between the onset of the annihilation or the scattering/cascade components and the blast wave synchrotron component.

In addition, the blast wave could also have a synchrotron self-Compton (SSC) component, from up-scattering of the synchrotron photons by the same electrons of comoving Lorentz factor γ' which produced them. The synchrotron peak photons x'_{sy} (and those in the power law above it) would scatter in the Klein-Nishina regime, since in the rest frame of the electrons $x'_{sy}\gamma'_m \geq 16$. Thus, the upscattered comoving photons would have a comoving frame peak energy $\sim \gamma'_m \sim 12 \ \text{GeV}\epsilon_{e,-1}\Gamma_{d,2.1}$, and in the observer frame these would appear at

$$E_{ssc}^{ob} \sim 75 \ \text{GeV} \ \epsilon_{e,-1}\Gamma_{d,2.1}^2 (20/[1+z]) \quad \text{(SSC)} \tag{11}$$

This SSC component would have a typical time lag relative to the peak synchrotron emission, lagging by

$$t_{ssc,lag}^{ob} \sim (r_d/c\Gamma_d^2)(1+z) \sim 4 \times 10^2 t_{a,3.5}([1+z]/20) \ \text{s} \ , \tag{12}$$

behind the synchrotron peak (10).

Reverse shock radiation?.– If the jet material, beyond some radius, becomes baryon-loaded one would expect a reverse shock. In this case, during the period when the external shock Lorentz factor decreases $\Gamma \propto r^{-1/2}$ before reaching the value Γ_d at the deceleration radius r_d (eq.[6]), the reverse shock Lorentz factor would increase as $\Gamma_r \propto r^{1/2}$ and could become relativistic. The reverse shock radiation would peak at r_d at the same time t_{peak}^{ob} as the forward shock synchrotron component, and it should have a comparable bolometric luminosity. To calculate the detailed properties of this reverse shock would, however, require knowledge about the baryon contamination level of the magnetic ejecta, which is highly speculative. Nonetheless, if the properties of the rare prompt optical flashes of lower redshift GRBs may be extrapolated to population III redshifts, we can get a speculative estimate. Taking as an example either GRB 080319B at $z = 0.937$ with $m_V \sim 5$ or GRB 050904 at $z = 6.29$ with $m_V \sim 10$, taking the ratios

of the squares of the luminosity distances one would expect an optical/IR flash of up to ~ 13 magnitudes, on similar timescales as the above X-ray and GeV components. On the other hand, if the baryon contamination was very low, a reverse shock would not form, since the Alfvénic sound speed would approach the speed of light (e.g. Giannios et al. (2008)).

Afterglow.– After the blast wave radiation has peaked, in an approximately uniform external density the Lorentz factor would decrease as $\Gamma \propto r^{-3/2}$ and the usual afterglow would set in, with a power law time decay of the flux and a softening of the spectrum. The properties of this afterglow would be fairly standard, except for any differences in the absorption properties of the metal-poor host environment.

5. Discussion

In this paper we have explored in some detail the spectral and temporal properties of high redshift population III GRBs within the context of a Poynting dominated relativistic jet model. The core collapse of a pop. III very massive star of $250 \lesssim (M_*/M_\odot) \lesssim 10^3$ will result in an intermediate mass black hole and a temporary accretion torus, and for fast rotating objects the extraction of rotational energy from the black hole could power a Poynting dominated outflow. The largest uncertainty, in such a model, is the value of the Poynting luminosity L and its time-dependence. The default assumption made here is that, for a constant magnetic α viscosity in the torus, L is approximately constant for $1/\alpha$ times the free-fall time from the boundary of the star. This is of course uncertain because the efficiency of field build-up is unknown, as is whether the magnetic stresses get up to a significant fraction of the gas pressure P. The quantity that determines the jet luminosity is the field strength around the hole, which depends on the peak density (and peak pressure P) near the inner part of the disk. In the standard alpha model the density in the disk goes as $r^{-3/2}$. However, for a radiation dominated gas (which is more compressible, even if the radiation is trapped on the relevant timescale) it would in principle be possible for the density law to be closer to r^{-3} (and this could happen if the effective α were to decrease towards small r). Thus, the disk could have a higher peak density (and steeper profile) whatever the initial stellar profile was – and the jet could have a much higher luminosity for a shorter period. Our timescale estimates are also subject to uncertainty. We took nominally the star to be rotating at half the break-up speed, but the stellar angular velocity could be non-uniform. If the outer regions were rotating more slowly than we assumed, the the disk would obviously be smaller, but the times scale would still be the free-fall time, so the accretion time could be up to a factor $\alpha = 10^{-1}\alpha_{-1}$ shorter. Thus, there are large uncertainties in the timescales and

luminosities we have estimated, which could be off in either direction. Nonetheless, the simplifying assumption made here may be appropriate considering the preliminary state of knowledge about population III stars.

The spectrum of the 'prompt' emission within the first day is shown to extend from soft X-rays to the multi-GeV range, with a characteristic time evolution. As a rough estimate, the luminosity of the annihilation photosphere (mainly in the form of photons) can be taken to be of the order of the remaining kinetic luminosity $L/2$, with an efficiency of conversion into annihilation photons of $\eta_\gamma \sim 1/2$. The annihilation component of eq.(7) peaks at 50 keV and extends as $N(E) \propto E^{-2}$ up to $\lesssim 3k m_e c^2$ in the comoving frame, or ~ 1 MeV in the observer frame. The the spectral energy flux per decade $E^2(dN/dE)$ is the total energy times $1/(\ln(E_{max}/E_{min}) = 1/\ln(20) \sim 0.31$. The radiation falling in the BAT band, 50-150 keV has $1/\ln(150/50) \sim 0.9$, so we can take the spectral efficiency in the BAT range, say 50-150 keV, as $\eta_{BAT} \sim 0.5 \times 0.3 \times 0.9 \sim 0.13 \sim 10^{-1}$. The X-ray flux might be even larger than this, adding a roughly comparable contribution from the external shock synchrotron component of eq. (9). With this efficiency, the predicted X-ray flux would be detectable in X-rays and hard X-rays by instruments such as the BAT detector on Swift, as estimated in §2. This would be detectable also in the GBM detector on Fermi, whose sensitivity is slightly better than Swift's. The GeV range Large Area detector (LAT) on Fermi has a fluence sensitivity for times $t \gtrsim 3 \times 10^4$ s of $\sim 3 \times 10^{-9} t^{1/2}$ even with significantly less than 10% of the luminosity in the GeV band at $t \sim 10^5$ s, this component would be detectable by the LAT. To detect them, however, it may be necessary to adjust the flux trigger algorithms to respond to a low level, very extended increase in the flux.

The spectral signature would have an initial hard, ~ 50 keV X-ray rise from the annihilation photons (equation [7]) lasting for about a day, with a possible extension out to ~ 25 GeV from up-scattering in the pair photosphere (equation [8]). There could be a cascade component from external photons leading almost simultaneously to soft X-rays in the few keV range, which is subject to considerable uncertainties. These would be followed, after a delay of hours up to a day, by an external shock synchrotron component in the few keV range (equation [9]). An inverse Compton component at energies in the 70 GeV range (equation [11]) may also be expected, lagging by about ten minutes after the keV range external shock synchrotron component. If the jets acquire a non-negligible baryon load at some stage before the external shock, a reverse shock may result in an infrared flash of $\gtrsim 13$th magnitude. An afterglow similar to that of lower redshift GRBs would follow over the next days, gradually shifting into the optical, infrared and radio frequency bands.

The detection of such very high redshift GRBs would be of great value, as it might

be the first and perhaps the only way to trace the formation of the first generation of stellar objects in the Universe. It could give important information about the redshift at which the initial objects form, the rate at which they form, and the input of radiation into the Universe at those early epochs and its contribution to the reionization of the intergalactic medium.

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